

# Custodial bulk Randall-Sundrum model and $B \rightarrow K^* l^+ l^-$

Sanghyeon Chang\* and C.S. Kim†

*Department of Physics and IPAP, Yonsei University, Seoul 120-749, Korea*

Jeonghyeon Song‡

*Department of Physics, Konkuk University, Seoul 143-701, Korea*

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## Abstract

The custodial Randall-Sundrum model based on  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$  generates new flavor-changing-neutral-current (FCNC) phenomena at tree level, mediated by Kaluza-Klein neutral gauge bosons. Based on two natural assumptions of universal 5D Yukawa couplings and no-cancellation in explaining the observed standard model fermion mixing matrices, we determine the bulk Dirac mass parameters. Phenomenological constraints from lepton-flavor-violations are also used to specify the model. From the comprehensive study of  $B \rightarrow K^* l^+ l^-$ , we found that only the  $B \rightarrow K^* e^+ e^-$  decay has sizable new physics effects. The zero value position of the forward-backward asymmetry in this model is also evaluated, with about 5% deviation from the SM result. Other effective observables are also suggested such as the ratio of two differential (or partially integrated) decay rates of  $B \rightarrow K^* e^+ e^-$  and  $B \rightarrow K^* \mu^+ \mu^-$ . For the first KK gauge boson mass of  $M_A^{(1)} = 2 - 4$  TeV, we can have about 10 – 20% deviation from the SM results.

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\*Electronic address: schang@cskim.yonsei.ac.kr

†Electronic address: kim@cskim.yonsei.ac.kr

‡Electronic address: jhsong@konkuk.ac.kr

## I. INTRODUCTION

The standard model (SM) has been very successful in reproducing nearly all experimental data on the fundamental interaction among gauge bosons and fermions. Nevertheless, the SM is not regarded as a fully satisfactory theory since it cannot explain two kinds of hierarchy: One is the gauge hierarchy, and the other is the hierarchy among the SM fermion masses. In the SM, the hierarchies are attributed to the hierarchical parameters – the bare Higgs mass parameter for the gauge hierarchy and the Yukawa couplings for the fermion mass hierarchy.

Randall and Sundrum (RS) scenario with bulk fermion fields is one of the rare candidates to be able to explain both hierarchies [1, 2, 3, 4]. The gauge hierarchy problem is explained by a geometrical exponential factor. Small SM fermion masses, which are proportional to the overlapping probability of the bulk fermion wave function with the confined Higgs boson field at the TeV brane, can be generated with moderate values of the bulk Dirac mass parameters [5, 6, 7, 8, 9, 10]. Yet the naive bulk RS model suffers from the strong constraints of the electroweak precision data (EWPD): The first Kaluza-Klein (KK) mode mass should be above  $\sim 20$  TeV [11, 12, 13, 14]. This is due to the lack of  $SU(2)$  custodial symmetry. In Ref. [15], an attractive model was proposed such that the custodial symmetry is induced from  $AdS_5/CFT$  feature of bulk gauge symmetry of  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ .

One of the interesting features of this model is new flavor-changing-neutral-current (FCNC) *at tree level*, mediated by KK gauge bosons [16]. This is due to the misalignment between the gauge eigenstates and the mass eigenstates: The five-dimensional (5D) Yukawa interaction is not generally flavor-diagonal. The fermion mass eigenstates of different generations can couple with KK modes of a neutral gauge boson, as one example is depicted in

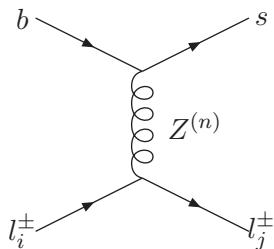


FIG. 1: Feynman diagram leading to  $B \rightarrow K^* l_i^+ l_j^-$  mixing in a warped extra dimension model.  $Z^{(n)}$  is the  $n$ -th KK mode of the SM  $Z$  boson.

Fig. 1. Note that FCNC at tree level in this model involves four external fermions: There is no tree-level effect on  $b \rightarrow s\gamma$ , for example. Since FCNC in the SM occurs only at loop level, rare FCNC decays can be a good place to probe the model.

Phenomenologically meaningful question is whether we have reliable predictions for various FCNC processes. The SM fermion mass matrix  $M_{ij}^f$  answers, which is determined by two ingredients. One is the 5D Yukawa couplings  $\lambda_{5ij}^f$ , and the other is the fermion mode function fixed by the bulk Dirac mass parameter  $c_F$ . Unfortunately, there is no unique way to determine both ingredients only from the observed SM fermion masses and mixing matrices, albeit extensive studies presenting the feasibility of the generation of SM fermion masses by controlling the bulk Dirac mass parameters.

One reasonable approach is to adopt minimal and natural assumptions. In this paper, we have two basic assumptions. The first one is that the 5D Yukawa couplings are universal, *i.e.*,  $\lambda_{5ij} \simeq \lambda_5 \sim \mathcal{O}(1)$ . Small masses of the SM fermions are explained by suppressed zero mode functions. Second, we assume that when explaining the observed SM mixing matrices, Cabibbo-Kobayashi-Maskawa (CKM) and Pontecorvo–Maki–Nakagawa–Sakata (PMNS) matrices, each of which is the product of two independent mixing matrices in this model, no order-changing by cancellation is allowed. Our choice has the least hierarchy, which is consistent with the main motivation of this model. Based on these two assumptions we examine whether all the bulk Dirac mass parameters as well as mixing matrices can be fixed, and whether we have reliable predictions for the phenomenological signatures of FCNC process such as  $B \rightarrow K^*l^+l^-$  [17, 18]. This is our primal goal.

In the quark sector, our two natural assumptions are to be shown enough to fix all the bulk Dirac mass parameters. In the lepton sector, there are some ambiguities due to the observed *large* mixing angles. We will examine the constraints from lepton-flavor-violating processes and determine the bulk lepton sector fairly accurately, which is one of our new results.

With the phenomenologically specified parameters, we will study the effect of the custodial bulk RS model on various observables of  $B \rightarrow K^*l^+l^-$ . This decay mode, especially with  $K^*$ , has several virtues in the experimental aspect. As well as producing very clean signature, its branching ratio is larger than the decay into  $K$ . In addition, a vector boson  $K^*$  decaying to  $K\pi$  allows us various angular analysis to measure many observables, such as the forward-backward asymmetry  $A_{FB}$ .  $A_{FB}$  is a very good observable to probe new

physics effect, since the so-called zero value position of  $A_{FB}(\hat{s}_0) = 0$  has strongly suppressed hadronic uncertainty in the calculation of the form factors. In addition, we present other sensitive probes of this model such as the ratio of differential decay rates for  $B \rightarrow K^* e^+ e^-$  and  $B \rightarrow K^* \mu^+ \mu^-$ . The sensitivity is due to sizable coupling of  $Z^{(1)}-e^+-e^-$  but suppressed coupling of  $Z^{(1)}-\mu^+-\mu^-$  in this model.

The organization of the paper is as follows. In Sec. II, we briefly review the custodial bulk RS model with  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ . In Sec. III, we formulate the bulk fermion sector, and determine all the bulk Dirac mass parameters based on our two natural assumptions. We will show the inevitable ambiguity in the lepton sector due to the large mixing angles. Section IV deals with the FCNC in this model. In Sec. V, we examine the lepton-flavor violating processes, and the new effect on  $B \rightarrow K^* l^+ l^-$ . Unique and sensitive observables to this model are also proposed. We conclude in Sec. VI.

## II. THE BULK RS MODEL: BASIC FORMULAE

The RS model is based on a 5D warped spacetime with the metric [1]

$$ds^2 = e^{-2\sigma(y)}(dt^2 - d\vec{x}^2) - dy^2, \quad (1)$$

where the fifth dimension  $y \in [0, L]$  is compactified on the  $S^1/\mathbb{Z}_2 \times \mathbb{Z}'_2$  orbifold, and the warped function is  $\sigma(y) = k|y|$  with  $k$  at the Planck scale  $M_{\text{Pl}}$ . There are two reflection symmetries under  $\mathbb{Z}_2 : y \rightarrow -y$  and  $\mathbb{Z}'_2 : y' (= y - L/2) \rightarrow -y'$ . Two boundaries are the  $\mathbb{Z}_2$ -fixed point at  $y = 0$  (Planck brane), and the  $\mathbb{Z}'_2$ -fixed point at  $y = L$  (TeV brane). In what follows, we denote  $(\mathbb{Z}_2, \mathbb{Z}'_2)$  parity by  $(\pm, \pm)$ . In many cases, conformal coordinate  $z \equiv e^{\sigma(y)}/k$  is more convenient:

$$ds^2 = \frac{1}{(kz)^2}(dt^2 - dx^2 - dz^2). \quad (2)$$

With  $kL \approx 35$ , the natural cut-off of the theory  $T \equiv e^{-kL}k$  becomes at the TeV scale, which answers the gauge hierarchy problem:

$$T \equiv \epsilon k \sim \text{TeV} \quad \text{with } \epsilon \equiv e^{-kL} \sim \frac{\text{TeV}}{M_{\text{Pl}}}. \quad (3)$$

We adopt the model suggested by Agashe *et.al.* in Ref. [15], based on the gauge structure of  $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ : The custodial symmetry is guaranteed by the bulk

$SU(2)_R$  gauge symmetry. The bulk gauge symmetry  $SU(2)_R$  is broken into  $U(1)_R$  by the orbifold boundary conditions on the Planck brane such that gauge fields  $\widetilde{W}_R^{1,2}$  have  $(-+)$  parity. The  $U(1)_R \times U(1)_{B-L}$  is spontaneously broken into  $U(1)_Y$  on the Planck brane, and the Higgs field localized on the TeV brane is responsible for the breaking of  $SU(2)_L \times U(1)_Y$  to  $U(1)_{EM}$ .

A 5D gauge field  $A^M(x, z)$  is expanded in terms of KK modes,

$$A_\nu(x, z) = \sqrt{k} \sum_n A_\nu^{(n)}(x) f_A^{(n)}(z), \quad (4)$$

where the zero mode function is  $f_A^{(0)} = 1/\sqrt{kL}$ . The massless zero mode is interpreted as a SM gauge field [3]. The general  $f_A^{(n)}(z)$  function can be found in early references, for example, in Ref. [8]. The bulk fermion field  $\Psi(x, z) \equiv e^{2\sigma} \hat{\Psi}$  is also expanded as

$$\hat{\Psi}(x, z) = \sqrt{k} \sum_n \left[ \psi_L^{(n)}(x) f_L^{(n)}(z) + \psi_R^{(n)}(x) f_R^{(n)}(z) \right]. \quad (5)$$

Two zero mode functions are

$$f_L^{(0)}(z, c) = f_R^{(0)}(z, -c) = \frac{(Tz)^{-c}}{N_L^{(0)}}, \quad (6)$$

where  $c$  is defined by the bulk Dirac mass  $m_D = ck \operatorname{sign}(y)$ , and  $N_L^{(0)}$  is referred to Ref. [8]. Note that a massless SM fermion corresponds to the zero mode with  $(++)$  parity. Since  $\Psi_L$  has always opposite parity of  $\Psi_R$ , the left-handed SM fermion is the zero mode of a 5D fermion whose left-handed part has  $(++)$  parity. The right-handed part has automatically  $(-)$  parity which cannot describe a SM fermion.

This characteristic feature of a bulk fermion in a warped model requires to extend the fermion sector. For each left-handed SM fermion, there should exist another bulk fermion whose right-handed part has  $(++)$  parity. Due to the gauge structure of  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ , these right-handed SM fermions belong to a  $SU(2)_R$  doublet. Since  $\widetilde{W}_R^{1,2}$  fields have  $(-+)$  parity and couple two elements of a  $SU(2)_R$  doublet, one of the  $SU(2)_R$  doublet should have  $(-+)$  parity. As a result, the whole quark sector is

$$Q_i = \begin{pmatrix} u_{iL}^{(++)} \\ d_{iL}^{(++)} \end{pmatrix}, \quad U_i = \begin{pmatrix} u_{iR}^{(++)} \\ D_{iR}^{(-+)} \end{pmatrix}, \quad D_i = \begin{pmatrix} U_{iR}^{(-+)} \\ d_{iR}^{(++)} \end{pmatrix}, \quad (7)$$

and lepton sector is

$$L_i = \begin{pmatrix} \nu_{iL}^{(++)} \\ e_{iL}^{(++)} \end{pmatrix}, \quad N_i = \begin{pmatrix} \nu_{iR}^{(++)} \\ E_{iR}^{(-+)} \end{pmatrix}, \quad E_i = \begin{pmatrix} N_{iR}^{(-+)} \\ e_{iR}^{(++)} \end{pmatrix}, \quad (8)$$

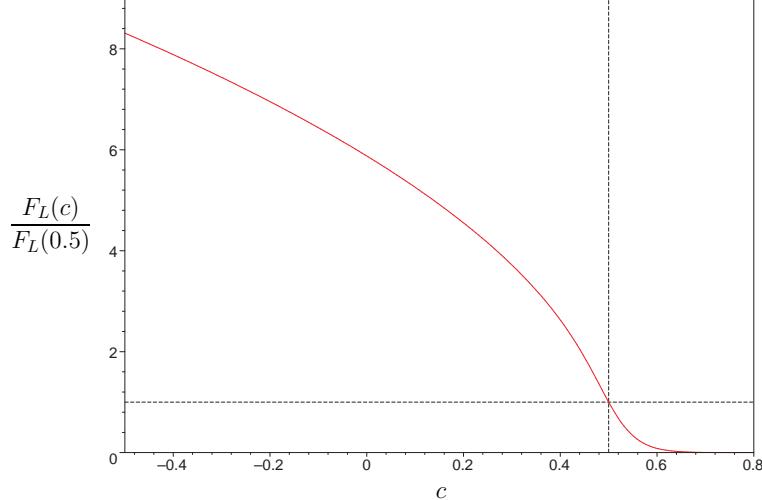


FIG. 2:  $F_L(c)/F_L(0.5)$  as a function of the bulk Dirac mass parameter  $c$ .

where  $i$  is the generation index. Dirac mass parameters  $(c_{Q_i}, c_{U_i}, c_{D_i}, c_{L_i}, c_{E_i}, c_{N_i})$  determine the fermion mode functions, KK mass spectra, and coupling strength with KK gauge bosons.

### III. THE SM FERMION MASSES AND MIXINGS

#### A. Basic Assumptions

On the TeV brane, the SM fermion mass is generated as the localized Higgs field develops its vacuum expectation value of  $\langle H \rangle = v \simeq 174$  GeV. The SM mass matrix for a fermion  $f (= u, d, \nu, e)$  is

$$(M_f)_{ij} = v \lambda_{5ij}^f \frac{k}{T} f_R^{(0)}(z, c_{R_i}) f_L^{(0)}(z, c_{L_j}) \Big|_{z=1/T} \equiv v \lambda_{5ij}^f F_R(c_{R_i}) F_L(c_{L_j}), \quad (9)$$

where  $i, j$  are the generation indices,  $\lambda_{5ij}^f$  are the 5D (dimensionless) Yukawa couplings and

$$F_L(c) = F_R(-c) \equiv \frac{f_L^{(0)}(1/T, c)}{\epsilon^{1/2}}. \quad (10)$$

Figure 2 shows  $F_L(c)$ , normalized by  $F_L(0.5)$ , as a function of  $c$ . The value of  $F_L(c)$  decreases with increasing  $c$ , and becomes suppressed once  $c > 0.5$ .

The mass eigenstates of the SM fermions are then

$$\chi_{fL} = U_{fL}^\dagger \psi_{fL}^{(0)}, \quad \chi_{fR} = U_{fR}^\dagger \psi_{fR}^{(0)}. \quad (11)$$

Note that the observed mixing matrix is a multiplication of *two independent* mixing matrices such that  $V_{\text{CKM}} = U_{uL}^\dagger U_{dL}$  and  $U_{\text{PMNS}} = U_{eL}^\dagger U_{\nu L}$ .

Due to the lack of *a priori* knowledge of bulk Dirac mass parameters and 5D Yukawa couplings, it is not generally possible to deduce all of their information only from the observed fermion mass spectrum and mixing angles. The number of unknown parameters far exceed the number of observations. One of the best approach is to develop the theory based on a few sound assumptions. We have the following two natural *assumptions*:

1. For all fermions, 5D Yukawa couplings have a common value  $\lambda_5$  of the order of one.
2. No order-changing by cancellation is allowed when the multiplication of two mixing matrices explains the observed mixing matrix.

For *assumption-1*, minor differences in  $\lambda_{5ij}$  are to be absorbed into mixing matrices. The top quark mass scale is naturally explained by  $v \simeq 174$  GeV. Other small SM fermion masses are generated by controlling  $c$ 's. The *assumption-2* is consistent with the spirit of no fine-tuning. When we write the elements of mixing matrices below, only their order of magnitude does matter.

The *assumption-1* leads to the following relation for the fermion mass matrix:

$$(M_f^T M_f)_{ij} = \lambda_5^2 v^2 F_L(c_{Li}) F_L(c_{Lj}) \sum_k F_R(c_{R_k})^2. \quad (12)$$

Since the left-handed up-type ( $u_{iL}$  or  $\nu_{iL}$ ) and down-type ( $d_{iL}$  or  $e_{iL}$ ) belong to the same  $SU(2)_L$  doublet and thus have the same  $c$ , Eq. (12) shows the proportionality of

$$M_u^T M_u \propto M_d^T M_d, \quad M_\nu^T M_\nu \propto M_e^T M_e. \quad (13)$$

Using the relation of

$$M_f^T M_f = U_{fL} (M_f^{(\text{d})})^2 U_{fL}^\dagger, \quad (14)$$

the proportionality in Eq. (13) helps determine the bulk Dirac mass parameters, if  $U_{fL}$  is known.

## B. Quark Sector Mass and Mixing

In the quark sector, *assumption-1* and -2 are enough to fix the model due to the hierarchical masses and almost diagonal mixing matrix. Nine Dirac mass parameters ( $c_{Q_i}$ ,  $c_{U_i}$ ,

$c_{D_i}$ ) are fairly well determined [7, 8, 9, 10]. The assumption-2 can be easily satisfied if both  $U_{uL}$  and  $U_{dL}$  are CKM-type: The  $V^{\text{CKM}} = U_{uL}^\dagger U_{dL}$  condition is naturally satisfied without any fine-tuned cancellation. We parameterize

$$(U_{qL})_{ij} = \kappa_{ij} V_{ij}^{\text{CKM}}, \quad (15)$$

where  $\kappa_{ij}$ 's are complex parameters of the order of one. To avoid order changing during the diagonalization of matrix, we take  $|\kappa_{ij}| \in [1/\sqrt{2}, \sqrt{2}]$ .

In the simplified Wolfenstein parametrization with  $\lambda \simeq 0.22$ , the CKM matrix is

$$V^{\text{CKM}} \simeq \begin{pmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}. \quad (16)$$

With the observed SM quark mass spectra of

$$M_u^{(\text{d})} \simeq v \text{ diag}(\lambda^8, \lambda^{3.5}, 1), \quad M_d^{(\text{d})} \simeq v \text{ diag}(\lambda^7, \lambda^5, \lambda^{2.5}), \quad (17)$$

we get

$$U_{uL}(M_u^{(\text{d})})^2 U_{uL}^T \simeq v^2 \begin{pmatrix} \lambda^6 & \lambda^5 & \lambda^3 \\ \lambda^5 & \lambda^4 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}, \quad U_{dL}(M_d^{(\text{d})})^2 U_{dL}^T \simeq v^2 \lambda^5 \begin{pmatrix} \lambda^6 & \lambda^5 & \lambda^3 \\ \lambda^5 & \lambda^4 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}. \quad (18)$$

Comparing two matrices in Eq. (18) based on Eqs. (13) and (14), we have

$$\begin{aligned} F_L(c_{Q_1}) : F_L(c_{Q_2}) : F_L(c_{Q_3}) &\simeq \lambda^3 : \lambda^2 : 1, \\ F_R(c_{A_1}) : F_R(c_{A_2}) : F_R(c_{A_3}) &\simeq \lambda^3 : \lambda^2 : 1, \quad \text{for } A = U, D, \\ F_R(c_{D_1})^2 + F_R(c_{D_2})^2 + F_R(c_{D_3})^2 &\simeq \lambda^5 [F_R(c_{U_1})^2 + F_R(c_{U_2})^2 + F_R(c_{U_3})^2]. \end{aligned} \quad (19)$$

Therefore, the SM quark mixing matrices can be approximated as

$$(U_{qL})_{ij(i \leq j)} \approx \frac{F_L(c_{Q_i})}{F_L(c_{Q_j})}, \quad (U_{qR})_{ij(i \leq j)} \approx \frac{F_R(c_{A_i})}{F_R(c_{A_j})}. \quad (20)$$

The bulk Dirac mass parameters are determined, as in Ref. [8],

$$\begin{aligned} c_{Q_1} &\simeq 0.61, \quad c_{Q_2} \simeq 0.56, \quad c_{Q_3} \simeq 0.3^{+0.02}_{-0.04}, \\ c_{D_1} &\simeq -0.66, \quad c_{D_2} \simeq -0.61, \quad c_{D_3} \simeq -0.56, \\ c_{U_1} &\simeq -0.71, \quad c_{U_2} \simeq -0.53, \quad 0 \lesssim c_{U_3} \lesssim 0.2. \end{aligned} \quad (21)$$

### C. Lepton Sector Mass and Mixing

The Pontecorvo–Maki–Nakagawa–Sakata (PMNS) mixing matrix in the weak charged lepton current is approximately

$$U_{\text{PMNS}} \simeq \begin{pmatrix} 0.8 & 0.5 & U_{e3} \\ 0.4 & 0.6 & 0.7 \\ 0.4 & 0.6 & 0.7 \end{pmatrix}, \quad (22)$$

where the current data constrains  $U_{e3} \lesssim 0.18$  at  $2\sigma$  [19]. Since  $U_{\text{PMNS}} = U_{eL}^\dagger U_{\nu L}$ , the elements of the (2, 3) block of  $U_{eL}$  and  $U_{\nu L}$  are of the order of one. In addition, the specific form of mass matrix in Eq. (12) allows only the normal mass hierarchy for the neutrino masses. The observed lepton masses are then

$$M_\nu^{(\text{d})} \simeq m_{\nu_3} \text{ diag}(0, \delta, 1), \quad M_e^{(\text{d})} \simeq m_\tau \text{ diag}(\delta^4, \delta^{1.5}, 1), \quad (23)$$

where  $\delta = \sqrt{\Delta m_{\text{sol}}^2 / \Delta m_{\text{atm}}^2} \approx 0.173$ . Numerical estimation shows that  $\delta \simeq U_{e3}$  in this model [8]. Then the condition of  $U_{eL}(M_e^{(\text{d})})^2 U_{eL}^\dagger \propto U_{\nu L}(M_\nu^{(\text{d})})^2 U_{\nu L}^\dagger$  from Eqs. (13) and (14) leads to

$$(U_{eL})_{13} \simeq (U_{\nu L})_{13}. \quad (24)$$

Using the unitarity condition of mixing matrices,  $U_{\nu L}$  is well constrained as

$$U_{\nu L} \sim \begin{pmatrix} 1 & 1 & \delta \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}. \quad (25)$$

Substituting  $U_{\nu L}$  in  $M_\nu^T M_\nu$ ,

$$M_\nu^T M_\nu = U_{\nu L}(M_\nu^{(\text{d})})^2 U_{\nu L}^\dagger \propto \begin{pmatrix} \delta^2 & \delta & \delta \\ \delta & 1 & 1 \\ \delta & 1 & 1 \end{pmatrix}, \quad (26)$$

we have the following relations among  $F_L(c_{L_i})$  from the definition in Eq. (12):

$$\delta \simeq \frac{F_L(c_{L_1})}{F_L(c_{L_2})} \simeq \frac{F_L(c_{L_1})}{F_L(c_{L_3})}. \quad (27)$$

From the behavior of  $F_L(c)$  in Fig. 2, we have the following hierarchy:

$$c_{L_2} \simeq c_{L_3} < c_{L_1}. \quad (28)$$

The relation  $U_{eL}(M_e^{(d)})^2 U_{eL}^\dagger \propto M_\nu^T M_\nu$  leaves minor ambiguity in  $U_{eL}$ :

$$U_{eL} \sim \begin{pmatrix} 1 & u_{12} & \delta \\ u_{21} & 1 & 1 \\ u_{31} & 1 & 1 \end{pmatrix}, \quad \text{for } u_{12} \lesssim \delta, \quad u_{21} + u_{31} \simeq \delta. \quad (29)$$

The condition of  $u_{21} + u_{31} \simeq \delta$  comes from the unitarity of  $U_{eL}$ , *i.e.*,  $(U_{eL}^\dagger U_{eL})_{13} = 0$ .

The matrix form of  $U_{eR}$  is well determined due to the hierarchical charged-lepton masses. We attribute  $F_R(c_{E_i})$  to the source of the hierarchy. The right-handed lepton mixing matrix should have an approximately symmetric form of

$$(U_{eR})_{ij} \approx \frac{F_R(c_{E_i})}{F_R(c_{E_j})} \quad \text{for } i \leq j. \quad (30)$$

Unlike in the quark sector, the lepton mass spectrum and mixing information are not enough to fix all the values of the bulk Dirac mass parameters for the  $SU(2)_L$  doublet. We will resort to the phenomenological constraint from the lepton flavor violating decays of  $\mu$  and  $\tau$  to reduce the ambiguity below.

#### IV. FCNC THROUGH KK GAUGE BOSONS

In this model, the mass eigenstate of the SM fermion is a mixture of gauge eigenstates as in Eq. (11). Since the 5D gauge interaction is flavor diagonal, we have FCNC mediated by KK gauge bosons. We denote  $W_L^{(n)}$ ,  $W_R^{(n)}$  and  $B_X^{(n)}$  for the KK gauge fields of  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ , respectively. Their 5D gauge couplings ( $g_{5L}$ ,  $g_{5R}$  and  $g_{5X}$ ) are related with the 4D effective couplings through

$$\begin{aligned} g &= g_{4L} = \frac{g_{5L}}{\sqrt{kL}}, \\ \tilde{g} &= g_{4R} = \frac{g_{5R}}{\sqrt{kL}} \sim g', \\ g_X &= g_{4X} = \frac{g_{4Y} g_{4R}}{\sqrt{g_{4R}^2 - g_{4Y}^2}} \sim g'. \end{aligned} \quad (31)$$

In terms of gauge eigenstates, the 4D gauge interactions with KK gauge modes are

$$\mathcal{L}_{4D} \supset g_{4D}^a \sum_{n=1}^{\infty} \left( \hat{g}_L^{(n)}(c_i) \bar{\psi}_{iL}^{(0)} T^a \gamma^\mu \psi_{iL}^{(0)} + \hat{g}_R^{(n)}(c_i) \bar{\psi}_{iR}^{(0)} T^a \gamma^\mu \psi_{iR}^{(0)} \right) A_\mu^{a(n)}, \quad (32)$$

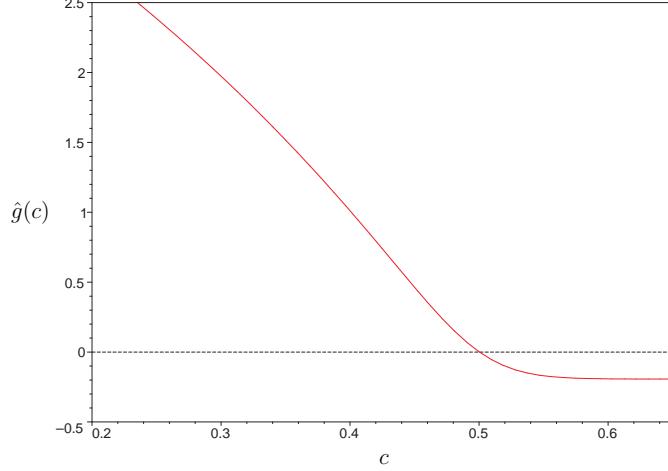


FIG. 3: Plot of  $\hat{g}^{(1)}(c)$  for bulk mass  $c$

where  $T^a = (T_L, T_R, Y_X)$  for  $A^a = (W_{3L}, W_{3R}, B_X)$ ,  $Y_X = (B - L)/2$ ,  $g_{4D}^a = g_5^a/\sqrt{kL}$  and

$$\begin{aligned}\hat{g}_L^{(n)}(c_{f_i}) &= \sqrt{kL} \int dz k \left[ f_L^{(0)}(z, c_{f_i}) \right]^2 f_A^{(n)}(z) \equiv \hat{g}^{(n)}(c_{f_i}), \\ \hat{g}_R^{(n)}(c_{f_i}) &= \sqrt{kL} \int dz k \left[ f_R^{(0)}(z, c_{f_i}) \right]^2 f_A^{(n)}(z) = \hat{g}^{(n)}(-c_{f_i}).\end{aligned}\quad (33)$$

For later discussions we plot  $\hat{g}^{(n)}(c)$  as a function of  $c$  in Fig. 3. Note that  $\hat{g}(c)$  vanishes at  $c = 1/2$ :

$$\hat{g}^{(n)} \left( c = \frac{1}{2} \right) = 0. \quad (34)$$

Another interesting feature is that the value of  $\hat{g}^{(n)}(c)$  converges into  $-0.2$  for  $c \gtrsim 0.55$ . As  $c$  becomes less than  $1/2$ , the value of  $\hat{g}^{(n)}(c)$  increases rather sharply.

Considering the  $B \rightarrow K^* l^+ l^-$  process, we focus on the mixing among the SM down-type quarks  $d_i$  and the SM charged leptons  $e_i$ , mediated by the  $n$ -th neutral KK gauge bosons in this model:

$$\begin{aligned}\mathcal{L}_{4D} \supset -\frac{1}{2} \sum_{i,j,n} \Bigg[ & g \left( K_{Qij}^{(n)} \bar{d}_{iL} \gamma^\mu d_{jL} + K_{Lij}^{(n)} \bar{e}_{iL} \gamma^\mu e_{jL} \right) W_{3L\mu}^{(n)} \\ & + \tilde{g} \left( K_{Dij}^{(n)} \bar{d}_{iR} \gamma^\mu d_{jR} + K_{Eij}^{(n)} \bar{e}_{iR} \gamma^\mu e_{jR} \right) W_{3R\mu}^{(n)} \\ & - g_X \left( K_{Qij}^{(n)} \bar{d}_{iL} \gamma^\mu d_{jL} - K_{Lij}^{(n)} \bar{e}_{iL} \gamma^\mu e_{jL} + K_{Dij}^{(n)} \bar{d}_{iR} \gamma^\mu d_{jR} - K_{Eij}^{(n)} \bar{e}_{iR} \gamma^\mu e_{jR} \right) B_{X\mu}^{(n)} \Bigg],\end{aligned}\quad (35)$$

where  $i, j$  are the generation indices ( $i, j = 1, 2, 3$ ), and

$$\begin{aligned}
K_{Qij}^{(n)} &= \sum_{k=1}^3 (U_{dL}^\dagger)_{ik} \hat{g}^{(n)}(c_{Q_k}) (U_{dL})_{kj} , \\
K_{Dij}^{(n)} &= \sum_{k=1}^3 (U_{dR}^\dagger)_{ik} \hat{g}^{(n)}(-c_{D_k}) (U_{dR})_{kj} , \\
K_{Lij}^{(n)} &= \sum_{k=1}^3 (U_{eL}^\dagger)_{ik} \hat{g}^{(n)}(c_{L_k}) (U_{eL})_{kj} , \\
K_{Eij}^{(n)} &= \sum_{k=1}^3 (U_{eR}^\dagger)_{ik} \hat{g}^{(n)}(-c_{E_k}) (U_{eR})_{kj} .
\end{aligned} \tag{36}$$

## V. FLAVOR VIOLATING PROCESS

### A. Lepton Flavor Violations

In this model, the flavor-violating interactions in Eq. (35) generates the lepton-flavor-violating decay of  $l \rightarrow l' l'' l'''$  at tree level, which is mediated by KK gauge bosons. Radiative lepton-violating processes such as  $\mu \rightarrow e\gamma$  does not happen at tree level. With negligible SM contributions, the bulk-RS effects become dominant for  $\tau \rightarrow 3e$  and  $\tau \rightarrow 3\mu$  with the following experimental bound [20]:

$$\frac{\Gamma(\tau \rightarrow 3e)}{\Gamma(\mu \rightarrow e\nu_\mu\bar{\nu}_e)} \simeq (K_{L11}^{(1)} K_{L12}^{(1)})^2 \left( \frac{m_Z}{M_A^{(1)}} \right)^4 \lesssim 1.0 \times 10^{-12}, \tag{37}$$

$$\frac{\Gamma(\tau \rightarrow 3\mu)}{\Gamma(\tau \rightarrow \mu\nu_\tau\bar{\nu}_\mu)} \simeq (K_{L22}^{(1)} K_{L23}^{(1)})^2 \left( \frac{m_Z}{M_A^{(1)}} \right)^4 \lesssim 10^{-6}, \tag{38}$$

where  $m_Z$  is the SM  $Z$  boson mass. Here we consider only the major contributions from the lightest KK gauge boson since the bulk-RS effect is suppressed by the forth power of  $M_A^{(1)}$ .

If  $M_A^{(1)} \leq 3$  TeV, Eqs. (37) and (38) constrain

$$c_{L2} \simeq c_{L3} \simeq 0.5. \tag{39}$$

We justify it for the case of  $M_A^{(1)} = 3$  TeV as follows. Substituting  $U_{eL}$  in Eq. (29) into  $K_{Lij}$  in Eq. (36), the  $\tau \rightarrow 3\mu$  constraint in Eq. (38) becomes

$$K_{L22}^{(1)} K_{L23}^{(1)} \simeq (\hat{g}_2 + \hat{g}_3)^2 \simeq (2\hat{g}_2)^2 < 1, \tag{40}$$

where  $\hat{g}_i \equiv \hat{g}^{(1)}(c_{L_i})$ , and the second equality comes from Eq. (27). We have also used  $\hat{g}_{1,2,3} \lesssim \mathcal{O}(1)$ , and  $u_{12,21,31} \lesssim \delta$ . From the functional behavior of  $\hat{g}(c)$  in Fig. 3, we have  $c_{L_2} \simeq c_{L_3} > 0.45$ . This mild condition on  $c_{L_2}$  and  $c_{L_3}$  constrains  $c_{L_1} > 0.55$  and thus  $|\hat{g}_1| \simeq 0.2$ , as can be seen from Eq. (27) and Figs. 2 and 3. More constraint on  $c_{L_{2,3}}$  comes from  $\mu \rightarrow 3e$ :

$$\begin{aligned} K_{L11}^{(1)} K_{L12}^{(1)} &\simeq \hat{g}_1 \{ \hat{g}_1 u_{12} + \hat{g}_2 (u_{21} + u_{31}) \} \\ &\simeq \hat{g}_1 \{ \hat{g}_1 u_{12} + \hat{g}_2 \delta \} \lesssim 10^{-3} \quad \text{for } M_A^{(1)} \simeq 3 \text{ TeV}, \end{aligned} \quad (41)$$

where the second equality is from Eq. (29). The saturating value  $|\hat{g}_1| \approx 0.2$  suppresses the  $u_{12}$  element of  $U_{eL}$  to be very small,  $u_{12} < \delta^2$ . In addition,  $\hat{g}_1 \hat{g}_2 \delta \lesssim 10^{-3}$  condition requires  $\hat{g}_2 < 0.05$ . It strictly constrains such that  $|c_{L_2} - 0.5| \leq 0.004$ . For  $M_A^{(1)} = 2$  TeV, the bound becomes even more strict:  $|c_{L_2} - 0.5| \leq 0.002$ .

As a natural solution for the lepton bulk mass parameters allowed by the current lepton-flavor violating processes, we choose

$$\begin{aligned} c_{L_1} &\simeq 0.59, \quad c_{L_2} \simeq 0.5, \quad c_{L_3} \simeq 0.5, \\ c_{E_1} &\simeq -0.74, \quad c_{E_2} \simeq -0.65, \quad c_{E_3} \simeq -0.55. \end{aligned} \quad (42)$$

## B. Effects on $B \rightarrow K^* l^+ l^-$

The FCNC decay  $B \rightarrow K^* l^+ l^-$  has been observed with the branching ratio of the order of  $10^{-6}$  [17], as well as the forward-backward asymmetry [18]. The total transition amplitude for  $b \rightarrow s l_i^+ l_j^-$  can be written as

$$\mathcal{M} = \mathcal{M}_{\text{SM}} + \mathcal{M}_{\text{new}}. \quad (43)$$

For the SM results, we refer to Ref. [21, 22]. For new physics contributions, we adopt the parametrization in Ref. [21],

$$\begin{aligned} \mathcal{M}_{\text{new}} = \frac{G_F \alpha}{\sqrt{2} \pi} V_{tb} V_{ts}^* \bigg[ &C_{LL} (\bar{s}_L \gamma^\mu b_L) (\bar{l}_L \gamma^\mu l_L) + C_{LR} (\bar{s}_L \gamma^\mu b_L) (\bar{l}_R \gamma^\mu l_R) \\ &+ C_{RL} (\bar{s}_R \gamma^\mu b_R) (\bar{l}_L \gamma^\mu l_L) + C_{RR} (\bar{s}_R \gamma^\mu b_R) (\bar{l}_R \gamma^\mu l_R) \bigg]. \end{aligned} \quad (44)$$

Note that other new physics parameters (*i.e.*,  $C_{LRLR}$ ) vanish in this model.

The RS contributions can be written as

$$\mathcal{M}_{\text{RS}} \simeq \sum_{n=1}^{\infty} \frac{1}{4M_A^{(n)2}} \left[ \begin{aligned} & \left( g^2 K_{Q23}^{(n)} K_{Lii}^{(n)} - g_X^2 K_{Q23}^{(n)} K_{Lii}^{(n)} \right) (\bar{s}_L \gamma^\mu b_L) (\bar{l}_{iL} \gamma^\mu l_{iL}) \\ & - g_X^2 K_{Q23}^{(n)} K_{Eii}^{(n)} (\bar{s}_L \gamma^\mu b_L) (\bar{l}_{iR} \gamma^\mu l_{iR}) \\ & - g_X^2 K_{D23}^{(n)} K_{Lii}^{(n)} (\bar{s}_R \gamma^\mu b_R) (\bar{l}_{iL} \gamma^\mu l_{iL}) \\ & + \left( \tilde{g}^2 K_{D23}^{(n)} K_{Eii}^{(n)} - g_X^2 K_{D23}^{(n)} K_{Eii}^{(n)} \right) (\bar{s}_R \gamma^\mu b_R) (\bar{l}_{iR} \gamma^\mu l_{iR}) \end{aligned} \right]. \quad (45)$$

Since physical observables are strongly suppressed by  $M_A^{(n)}$ , we consider only the first KK mode effect and we omit the KK mode number notation ( $n$ ) in the rest of this section.

The preferred  $c_Q$ 's in Eq. (21) and the CKM-type matrices  $U_{qL}$  and  $U_{qR}$  in Eq. (20) simplify  $K_{Q23}$  and  $K_{D23}$  as

$$K_{Q23} \simeq (U_{qL})_{23} (U_{qL})_{33} \hat{g}(c_{Q_3}) \equiv \kappa_Q^2 \hat{g}(c_{Q_3}) V_{tb} V_{ts}^*, \quad (46)$$

$$\begin{aligned} K_{D23} & \simeq \left[ (U_{dR})_{22} (U_{dR})_{32} \hat{g}(c_{D_2}) + (U_{dR})_{23} (U_{dR})_{33} \hat{g}(c_{D_3}) \right] \\ & \equiv 2 \kappa_D^2 \hat{g}(c_{D_3}) V_{tb} V_{ts}^*, \end{aligned} \quad (47)$$

where we have used  $\hat{g}(c_{Q_3}) \gg \hat{g}(c_{Q_{1,2}})$  and  $\hat{g}(c_{D_2}) \approx \hat{g}(c_{D_3})$ . New physics parameters  $C_{XX'}$  ( $X, X' = L, R$ ) are

$$\begin{aligned} C_{LL} & \simeq \left( \frac{\tilde{G}}{M_A} \right)^2 (g^2 - g_X^2) \kappa_Q^2 \hat{g}(c_{Q_3}) K_{Lij}, \\ C_{LR} & \simeq \left( \frac{\tilde{G}}{M_A} \right)^2 g_X^2 \kappa_Q^2 \hat{g}(c_{Q_3}) K_{Eij}, \\ C_{RL} & \simeq 2 \left( \frac{\tilde{G}}{M_A} \right)^2 g_X^2 \kappa_D^2 \hat{g}(c_{D_3}) K_{Lij}, \\ C_{RR} & \simeq 2 \left( \frac{\tilde{G}}{M_A} \right)^2 (\tilde{g}^2 - g_X^2) \kappa_D^2 \hat{g}(c_{D_3}) K_{Eij}, \end{aligned} \quad (48)$$

where  $\tilde{G} = (\pi/2\sqrt{2}G_F\alpha)^{1/2} \approx 3.5$  TeV.

In Table I, we present the values of  $C_{LL}$ ,  $C_{RL}$ ,  $C_{LR}$ , and  $C_{RR}$  for  $b \rightarrow sl_i^+ l_j^-$ . For representative purpose, we set  $M_A^{(1)} = 2$  TeV,  $\kappa_{Q,D} = 1$ ,  $\delta = 0.15$ , and use central values of  $c$ 's in Eqs. (21) and (42). The values of  $C_{XX'}$  can be understood from  $\hat{g}$ :

$$\hat{g}(c_{Q_3}) \simeq 2.0, \quad \hat{g}(c_{L_2}) \simeq \hat{g}(c_{L_3}) = 0, \quad (49)$$

$$\hat{g}(c_{D_3}) \simeq \hat{g}(c_{L_1}) \simeq \hat{g}(-c_{E_1}) \simeq \hat{g}(-c_{E_2}) \simeq \hat{g}(-c_{E_3}) \simeq -0.2.$$

TABLE I: The values of  $C_{LL}$ ,  $C_{RL}$ ,  $C_{LR}$ , and  $C_{RR}$  for  $b \rightarrow s l_i^+ l_j^-$ . We set  $M_A^{(1)} = 2$  TeV,  $\kappa_{Q,D} = 1$ , and  $\delta = 0.15$ .

	$e^+ e^-$	$e^+ \mu^-$	$e^+ \tau^-$	$\mu^+ \mu^-$	$\mu^+ \tau^-$	$\tau^+ \tau^-$
$C_{LL}$	-0.3	$\pm 7 \times 10^{-3}$	$\pm 0.05$	$-7 \times 10^{-3}$	$\pm 6 \times 10^{-3}$	$-4 \times 10^{-5}$
$C_{RL}$	0.02	$\pm 5 \times 10^{-4}$	$\pm 4 \times 10^{-3}$	$6 \times 10^{-4}$	$\pm 5 \times 10^{-4}$	$3 \times 10^{-6}$
$C_{LR}$	-0.1	$\pm 0.01$	$\pm 10^{-3}$	-0.1	$\pm 0.01$	-0.1
$C_{RR}$	0.03	$\pm 3 \times 10^{-3}$	$\pm 2 \times 10^{-4}$	0.03	$\pm 3 \times 10^{-3}$	0.02

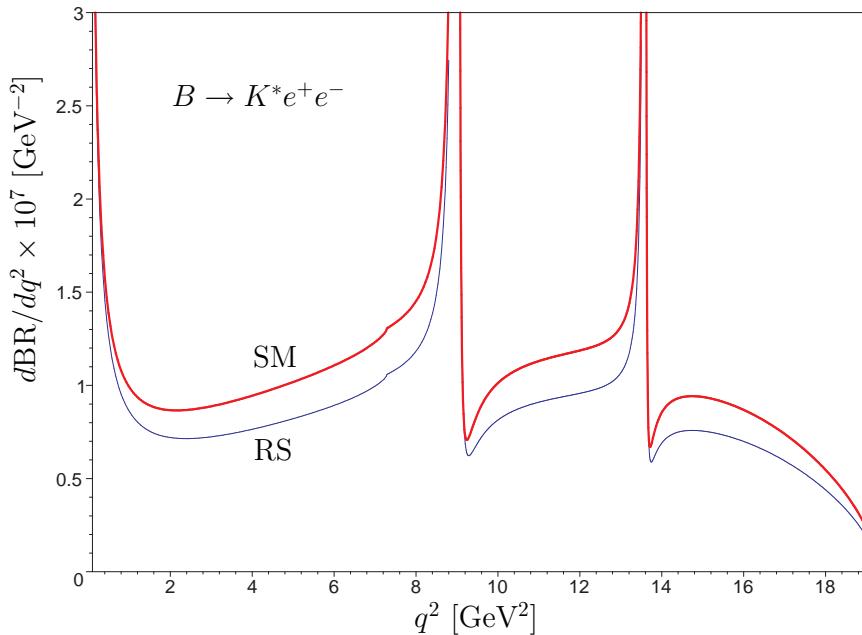


FIG. 4:  $dBR/dq^2$  as a function of  $q^2$  for  $B \rightarrow K^* e^+ e^-$ . The thick (red) line is the SM result, and the thin (blue) line is for the bulk RS model with  $\kappa = \sqrt{2}$  and  $M_A^{(1)} = 2$  TeV.

Brief comments on the sign of  $C_{XX'}$  are in order here. The negative signs of  $C_{LL}$  and  $C_{LR}$  are due to positive  $\hat{g}(c_{Q_3})$ , and negative  $\hat{g}(c_{L_1})$  and  $\hat{g}(-c_{E_1})$  which dominantly contribute to  $K_{Lij}$  and  $K_{Eij}$ , respectively. The sign of  $C_{XX'}$  for off-diagonal decays such as  $B \rightarrow K^* l^+ l^-$  is not determined since we could fix only the magnitude of elements of mixing matrices. In the magnitudes, only the  $C_{XX'}$ 's for  $b \rightarrow s e^+ e^-$  are substantial.  $C_{XX'}$ 's for decays involving  $\mu^\pm$  or  $\tau^\pm$  are quite suppressed, since  $\hat{g}(c_{L_2}) \simeq \hat{g}(c_{L_3}) \ll 1$ . Among  $C_{XX'}$ 's for  $b \rightarrow s e^+ e^-$ ,  $C_{LX}$  is larger than  $C_{RX}$  since  $\hat{g}(c_{Q_3})$  is much larger than  $\hat{g}(c_{D_3})$ .

In Fig. 4, we present the differential branching ratio  $dBR/dq^2$  as a function of  $q^2$  for

$B \rightarrow K^* e^+ e^-$ . We use the following values for the Wilson coefficients of the SM:

$$C_9^{\text{NDR}} = 4.153, \quad C_{10} = -4.546, \quad C_7 = -0.311, \quad (50)$$

which correspond to the next-to-leading QCD corrections [23, 24]. The renormalization scale  $\mu$  and the top quark mass are set to be

$$\mu = m_b = 4.8 \text{ GeV}, \quad m_t = 175 \text{ GeV}. \quad (51)$$

We follow Refs. [25] in taking into account the long-distance effects of the charmonium states. For the form factors, we have used the light-cone QCD sum-rule method predictions [26]. Throughout numerical analysis, we used the central values of the input parameters, and do not consider the theoretical uncertainty in the calculation of form factors. In Fig. 4, the thick (red) line is the SM result, and the thin (blue) line is for the bulk RS model. We have used the allowed maximum value of  $C_{XX}$ 's with  $\kappa (= \kappa_Q = \kappa_D) = \sqrt{2}$  and  $M_A^{(1)} = 2 \text{ TeV}$ . As discussed in Ref. [21], this BR distribution is most sensitive to  $C_{LL}$ . Since our  $C_{LL}$  for  $B \rightarrow K^* e^+ e^-$  is negative, the result in this model is less than in the SM. The reduction can be maximally about 20% at some points. Unfortunately, the theoretical uncertainty of the form factors are known to be about 15% [26]. It would be quite challenging for experiments to probe this new physics effect from the BR distribution.

One sensitive observable to new physics is known to be the zero value position of the forward-backward asymmetry, *i.e.*,  $A_{\text{FB}}(\hat{s}_0) = 0$ . The forward-backward asymmetry  $A_{\text{FB}}(\hat{s})$  is defined by

$$\frac{d}{d\hat{s}} A_{\text{FB}}(\hat{s}) = \frac{\int_0^1 dz \frac{d\Gamma}{d\hat{s} dz} - \int_{-1}^0 \frac{d\Gamma}{d\hat{s} dz}}{\int_0^1 dz \frac{d\Gamma}{d\hat{s} dz} + \int_{-1}^0 \frac{d\Gamma}{d\hat{s} dz}}, \quad (52)$$

where  $\hat{s} = q^2/m_B^2$ ,  $z = \cos \theta$ , and  $\theta$  is the angle between  $K^*$  and  $l^-$ . In the large energy expansion theory, it has been shown that  $\hat{s}_0$  has no hadronic uncertainty; it is determined simply by the short-distance Wilson coefficients  $C_9^{\text{eff}}$  and  $C_7^{\text{eff}}$  [27]. In Fig. 5, we show the  $A_{\text{FB}}(\hat{s})$  as a function of  $\hat{s}$ . The thick (red) line is the SM result, and the thin (blue) line is the new physics result with  $\kappa = \sqrt{2}$  and  $M_A^{(1)} = 2 \text{ TeV}$ . The zero value position of  $A_{\text{FB}}$  in the SM model is consistent with other results [28]. In our new model,  $\hat{s}_0$  shifts to the positive direction:  $\hat{s}_0$  can increase maximally about 18%. Experimental sensitivity is expected to reach this difference in near future.

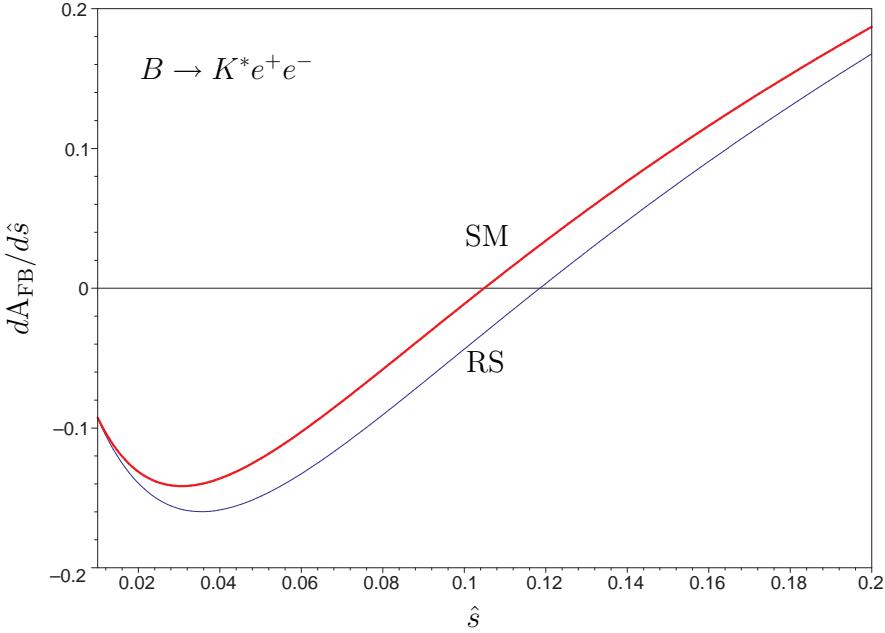


FIG. 5:  $dA_{FB}/d\hat{s}$  as a function of  $\hat{s}$  for  $B \rightarrow K^*e^+e^-$ . The thick (red) line is the SM result, and the thin (blue) line is the new physics result with  $\kappa = \sqrt{2}$  and  $M_A^{(1)} = 2$  TeV.

Now we present new phenomenological signatures exclusively for this model. One of the most unique features is that only the  $B \rightarrow K^*e^+e^-$  decay has sizable new physics effect while others have negligible effects. Therefore, we consider the ratio of differential decay rate of  $B \rightarrow K^*e^+e^-$  to that of  $B \rightarrow K^*\mu^+\mu^-$ . This ratio was proposed as an efficient observable to test the SM [29]. In Fig. 6, we show the ratios as a function of  $\hat{s}$  in the SM and the bulk RS model. The thick (red) line is for the SM result, the thin (black) line for the RS result with  $\kappa = 1$ , and the normal (blue) line for the RS result with  $\kappa = \sqrt{2}$ . In the most range of  $\hat{s}$ , the RS result is far below the SM one. For maximally allowed value of  $C_{XX'}$  with  $\kappa = \sqrt{2}$ , the deviation from the SM result can be about 20% for sizable range of  $\hat{s}$ . Even for moderate values of  $C_{XX'}$  with  $\kappa = 1$ , the deviation reaches up to 7%. Moreover, as taking the *ratio* of differential decay rates, most of the hadronic uncertainty in the calculation of form factors disappears. This can be a quite clean signal for experiments.

In order to see the dependence of new physics effect on  $M_A^{(1)}$ , we present the ratio of two partially integrated decay rates for  $B \rightarrow K^*e^+e^-$  and  $B \rightarrow K^*\mu^+\mu^-$ . From the profile in Fig. 7 as a function of  $\hat{s}$ , we choose the integration range of  $\hat{s} \in [0.1, 4m_c^2/m_B^2]$  with  $m_c$  being the charm quark mass. The dotted line is for the SM result, the dashed line for the RS result with  $\kappa = 1$ , and the solid line for the RS result with  $\kappa = \sqrt{2}$ . If  $C_{XX'}$ 's have allowed

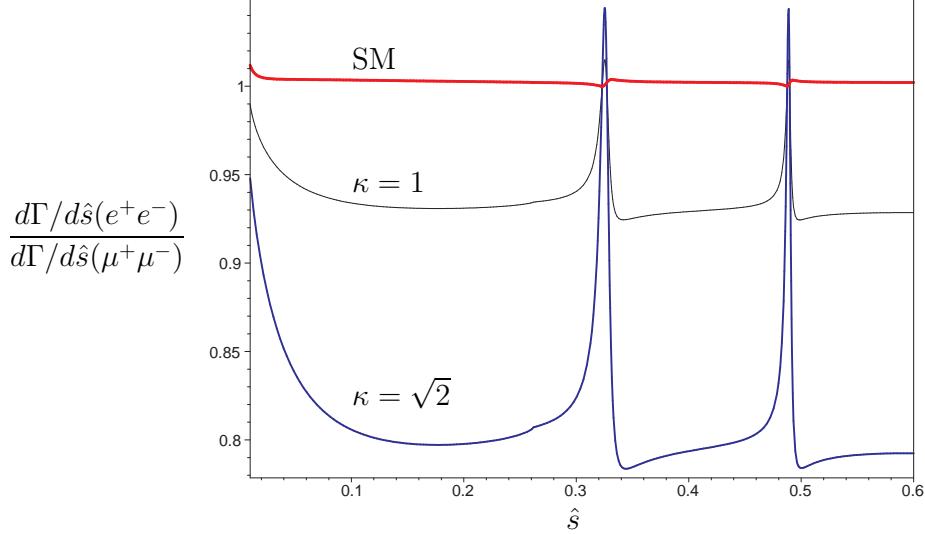


FIG. 6:  $\frac{d\Gamma}{d\hat{s}}(B \rightarrow K^* e^+ e^-) / \frac{d\Gamma}{d\hat{s}}(B \rightarrow K^* \mu^+ \mu^-)$  as a function of  $\hat{s}$  in the SM and the custodial bulk RS with  $\kappa = 1, \sqrt{2}$ . we set  $M_A^{(1)} = 2$  TeV.

maximum values ( $\kappa = \sqrt{2}$ ), the RS result with  $M_A^{(1)} = 2$  TeV shows about 18% deviation from the SM result, and that even with  $M_A^{(1)} = 4$  TeV shows about 4.5%. If  $C_{XX'}$ 's have medium values ( $\kappa = 1$ ), the RS result with  $M_A^{(1)} = 2$  TeV shows about 6% deviation from the SM result, and that with  $M_A^{(1)} = 4$  TeV shows about 2%. Since the ratio does not suffer from the hadronic uncertainty of the form factors, this difference will be within experimental sensitivity in near future.

## VI. CONCLUSIONS

The custodial Randall-Sundrum model is a warped 5D model with all the SM fields in the bulk. Only the Higgs boson field is confined on the TeV brane, which generates masses for the SM particles. The troublesome EWPD constraint is overcome by  $SU(2)$  custodial symmetry induced from  $AdS_5/CFT$  feature of bulk gauge symmetry of  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ . We focused on new FCNC phenomena which occur due to the misalignment between the gauge couplings and the 5D Yukawa interaction. We have the vertex of  $f-f'-A^{(n)}$ , where  $f^{(i)}$  is a SM fermion and  $A^{(n)}$  is a Kaluza-Klein mode of a neutral gauge boson. At tree level, we have non-SM FCNC involving four external SM fermions, mediated by KK neutral gauge bosons.

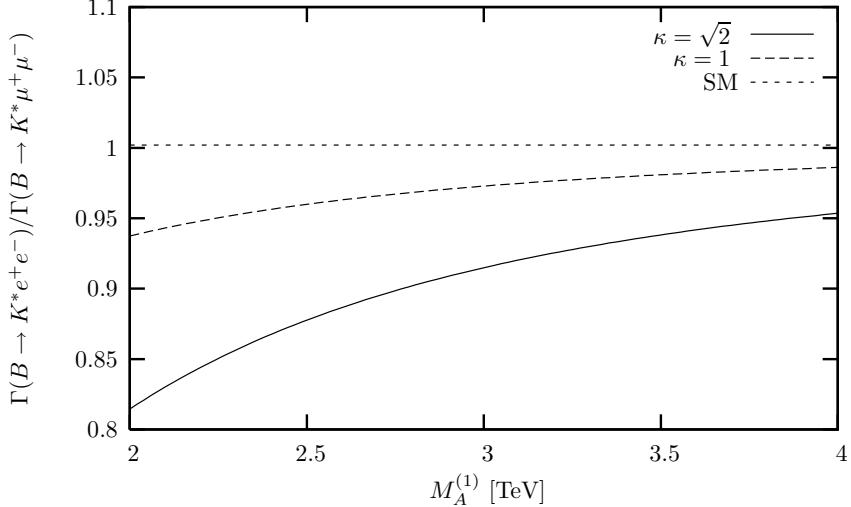


FIG. 7:  $\frac{\Gamma(B \rightarrow K^* e^+ e^-)}{\Gamma(B \rightarrow K^* \mu^+ \mu^-)}$  as a function of  $M_A^{(1)}$  in the SM and the custodial bulk RS with  $\kappa = 1, \sqrt{2}$ . Here  $\Gamma$  is partially integrated  $d\Gamma/d\hat{s}$  for  $\hat{s} \in [0.1, 4m_c^2/m_B^2]$ .

The  $f$ - $f'$ - $A^{(n)}$  vertex depends on two kinds of model parameters, the 5D Yukawa couplings and the bulk Dirac mass parameters. They also determine the SM fermion mass spectrum and mixing angles. Based on two natural assumptions of universal 5D Yukawa coupling and no-cancellation in explaining the observed SM fermion mixing matrices, we have obtained all the information on  $c$ 's as well as mixing angles.

In the custodial bulk RS model with very specified fermion structure, we study FCNC process of  $B \rightarrow K^* l^+ l^-$ . New physics effect is parameterized in the helicity amplitude as  $C_{XX'}(\bar{s}_X \gamma^\mu b_X)(\bar{l}_{X'} \gamma^\mu l_{X'})$ , where  $X, X' = L, R$ . If  $C_{XX'}$ 's have maximally allowed values, the differential decay rate of  $B \rightarrow K^* e^+ e^-$  deviates from the SM result as much as about 20% at some  $q^2$ . Unfortunately, the hadronic uncertainty in the form factors is large enough to sweep away this new effect. Instead, the zero value point of the forward-backward asymmetry,  $\hat{s}_0$ , is known to be quite insensitive to the hadronic uncertainty. In the maximal case, the deviation of  $\hat{s}_0$  from the SM value is about 5%, which is expected to be probed in near future.

We have also found the following characteristic features:

- The best chance to observe the custodial bulk RS model effect is through  $b \rightarrow s e^+ e^-$  due to the suppressed couplings of  $\mu^+ \mu^- Z^{(n)}$  and  $\tau^+ \tau^- Z^{(n)}$ . And  $C_{LL}$  is dominant, and  $C_{LR}$  is the second dominant.

- Two other decays of  $b \rightarrow s\mu^+\mu^-$  and  $b \rightarrow s\tau^+\tau^-$  have dominant vertex of  $C_{LR}$ . Unfortunately, their magnitudes are too small for experiments to probe in near future.
- Other non-diagonal decay modes of  $b \rightarrow sl_i^+l_j^- (i \neq j)$  are quite suppressed in this model.

Based on these observations, we suggested new phenomenological signatures to probe the custodial bulk-RS model. The first one is the ratio of two differential decay rates of  $B \rightarrow K^*e^+e^-$  and  $B \rightarrow K^*\mu^+\mu^-$ . Upon taking the ratio, the hadronic uncertainty in the calculation of form factors becomes negligible. Since two decay modes in the SM have almost the same decay rates with slight kinematic difference from the lepton masses, dominant new physics effects only for  $B \rightarrow K^*e^+e^-$  leads to sizable deviation from the SM result. For  $M_A^{(1)} = 2$  (4) TeV, the deviation can reach about 20% (7%). We also showed the ratio of partially integrated decay rates, which shows also about 10-20% deviation from the SM results. This deviation is expected to be observed in near future.

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